

Explosive Collisions at RHIC ?

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Motivated by experimental results from RHIC, we suggest how a condensate for the Polyakov loop might produce explosive behavior at the QCD phase transition. This is due to a rapid rollover of the condensate field below the transition temperature.

Many new results have poured forth from RHIC ($\sqrt{s} = 130$ A GeV). Several appear to be qualitatively different from those at the SPS ($\sqrt{s} = 17$ A GeV). “Dynamical” fluctuations in the average transverse momentum, p_t , are almost three times larger than at the SPS [1]. Observed HBT radii are small, $\sim 5 - 7$ fm [2], and average transverse boost velocities are large, $\sim .6$ c [3]. In particular, that the HBT radii $R_{out}/R_{side} \leq 1$, as a function of pion pair $p_t : 200 \rightarrow 400$ MeV, suggests “explosive” behavior [2]. Lastly, suppression of particles at high $p_t : 2 \rightarrow 6$ GeV appears to be a clear diagnostic probe [4,5].

Based upon a condensate model [6], previously we predicted dynamical fluctuations in the average transverse momentum [7]. Using known features of the condensate model [7], here we suggest how it might produce explosive behavior at the QCD phase transition, and affect the suppression of particles at high p_t .

At infinitely high temperature, by asymptotic freedom QCD is an ideal gas of massless quarks and gluons. Lattice data [8] and effective theories [9] suggest that from infinite temperature, on down to perhaps twice the transition temperature, T_c , it really is a plasma: a nearly ideal gas of quarks and gluons, augmented with “thermal” masses.

For temperatures of $2T_c$ down to T_c , instead of quasiparticles, it is more useful to think of a condensate [6]. In a nonabelian gauge theory without quarks, the order parameter for the deconfining phase transition is the Polyakov loop, ℓ . This is the trace of the $SU(3)$ color Aharanov-Bohm phase factor in the imaginary time direction. By asymptotic freedom, the expectation value of ℓ , $\langle \ell \rangle = \ell_0$, is one at infinite temperature; from ’t Hooft, it is zero in the confined phase, $T < T_c$. Thus in going from $2T_c$ down to T_c , the condensate for the Polyakov loop evaporates.

In the condensate model, the pressure is determined by a *potential* for ℓ [6], with

$$p \sim \ell_0^4 p_{\text{ideal}} \quad , \quad p_{\text{ideal}} \sim T^4 \quad , \quad (1)$$

where p_{ideal} is the pressure for an ideal gas of quarks and gluons. In the pure gauge theory, where the quarks are quenched, the lattice finds that at $2T_c$, $p/p_{\text{ideal}} \sim .8$ [8], so maybe $\ell_0 \sim .95$.

Lattice data for p/p_{ideal} can then be used to fit $\ell_0(T)$. If e is the energy density,

$$\frac{e - 3p}{T^4} \sim T \frac{\partial \ell_0^4}{\partial T} = 4T \ell_0^3 \frac{\partial \ell_0}{\partial T}. \quad (2)$$

In the pure gauge theory, the lattice finds that there is a sharp “bump” in $(e - 3p)/T^4$ above T_c . This bump indicates that the transition occurs in a relatively narrow region of temperature. In the condensate model, this reflects a rapid change in the potential for ℓ_0 .

Although the lattice data is less reliable with dynamical quarks, it suggests that the pressure is “flavor independent” [8]. In a mean field analysis of the condensate model, this implies that the ℓ -potential is determined mainly by the gluons — for which there is good lattice data. Without quarks, the deconfining transition is weakly first order. We make the most conservative assumption, that quarks wash out the deconfining transition, to leave only crossover. Even so, with dynamical quarks, a bump in $(e - 3p)/T^4$ is still evident [8], indicating a rapid change in the ℓ -potential near T_c .

For heavy ion collisions at high energies, the central region exhibits rapid longitudinal expansion along the beam direction. For an ideal gas, with boost invariant Bjorken expansion the temperature falls as $T \sim 1/\tau^{1/3}$, where τ is the proper time.

If the transition were strongly first order (as occurs for four or more colors), a mixed phase exists, and a hydrodynamic analysis indicates that the system lasts for a long time at T_c [10]. Because of this, the HBT ratio $R_{\text{out}}/R_{\text{side}}$ is significantly larger than one, ~ 2 [11]. Thinking of the firetube in the central region as a log, this is the “burning log” scenario [11]. A strong first order transition implies that there is a mixed phase of bubbles at T_c . For slow expansion, bubbles with $\ell_0 = 0$ coexist, at equal pressure, with those with $\ell_0 \neq 0$ [12].

If the transition is crossover, and the potential for ℓ changes slowly, then the system evolves smoothly through T_c , with a gradual shift of ℓ_0 from ~ 1 to ~ 0 . This is a “smoldering log”, with $R_{\text{out}}/R_{\text{side}} > 1$.

If the transition is crossover (or weakly first order), *and* the potential for ℓ changes *rapidly*, then the system quickly passes through T_c . Suddenly, the system finds itself below T_c , at a nonzero value of ℓ_0 , which is no longer the minimum. It then rolls down, oscillating about zero, until it settles to $\ell_0 = 0$. How this roll down occurs is described by the following equation:

$$Z_0 \frac{\partial^2}{\partial \tau^2} \ell + \gamma_\ell \frac{\partial}{\partial \tau} \ell - Z_s \nabla^2 \ell + \mathcal{V}'(\ell) + h_\Phi \Phi^2 \ell = 0. \quad (3)$$

Here ∇^2 is the Laplacian in the spatially transverse and rapidity directions, and $\mathcal{V}(\ell)$ is the potential for ℓ . The Polyakov loop is coupled to a field Φ for Goldstone bosons (π , K , η , η') through a term $\sim h_\Phi \Phi^2 \ell$ [7]. In the condensate model, the potential $\mathcal{V}(\ell)$ is completely fixed by the pressure. The coefficient for spatial variations in the Polyakov loop, Z_s , was taken from that for $SU(3)$ Wilson lines, $Z_s \sim 1/g^2$ [7]; actually, $Z_s \sim g^4$. We also assumed a Lorentz invariant form, $Z_0 = Z_s$, although probably $Z_0 \neq Z_s$.

In [7] we analyzed this equation in the Hartree approximation, where $\Phi^2 \rightarrow \langle \Phi^2 \rangle$, which neglects collisions amongst the produced pions. The condensate field ℓ was assumed to be spatially homogeneous, so that only its variation in time mattered. We took $\gamma_\ell = 0$; this neglects the effects of longitudinal expansion, which contributes $1/\tau$ to γ_ℓ , and dissipation.

Above T_c , the potential is sharp, and ℓ is trapped in a minimum with $\ell_0 \neq 0$. For a window of temperature which is extremely narrow, $\Delta T \approx 2\%T_c$, the potential changes suddenly into that for the symmetric phase: ℓ rolls down, and oscillates about zero. The initial energy density, which is stored entirely in the ℓ field, is then pumped into the production of pions. As the initial energy density is large, the produced pions have large average transverse momentum, $\sim 5f_\pi = 500$ MeV. We note that while this agrees with the experimentally measured result [3], in the present model this value cannot be tuned. The system is underdamped, and the average field oscillates several times about zero. These oscillations produce relatively large fluctuations about the average transverse momentum, p_t [7], on the order of $\sim 10\%$. This scenario is like reheating after inflation, or a quench in Disoriented Chiral Condensates (DCC's) [7,13]. It appears unlikely, however, that many coherent oscillations of an ℓ -field, sloshing back and forth, generates explosive behavior.

The STAR results suggest an alternate scenario. Perhaps, due either to rapid longitudinal expansion, or to an intrinsic term $\gamma_\ell \neq 0$, one is in a regime where the effects of damping cannot be neglected. In this case, the potential evolves quickly, but because of the damping, ℓ_0 oscillates only a few times. It is conceivable — but not guaranteed — that in this case particle production is explosive. Particles are emitted from a shell [14], with $R_{out}/R_{side} \leq 1$.

A semiclassical calculation of particle production in the condensate model appears reasonable. For a large number of colors, N_c , the potential $\mathcal{V}(\ell) \sim N_c^2$ [6], while Z_0 , Z_s , $\gamma_\ell \sim 1$; thus a semiclassical calculation is nominally valid to $\sim 1/N_c^2 \sim 10\%$ for $N_c = 3$. Getting spectra which are exponential in transverse mass is probably easy [15]; getting relative ratios right, such as $K/\pi \sim 0.2$, is far from obvious. Explosive behavior should also produce a characteristic electromagnetic spectrum, such as in direct photons [16].

Both scenarios produce relatively large fluctuations about the average transverse momentum, which are measurable on an Event-by-Event basis [17]. In the underdamped scenario, production is from causally connected regions, of size ~ 1 fm [7]. Due to Poisson statistics, the fluctuations from each domain are then reduced by the inverse square root of the number of domains in one unit of rapidity; If in one unit of rapidity there are ~ 300 domains, the fluctuations are $\sim 10\%/\sqrt{300} \sim .6\%$. Further, smaller rapidity bins should increase the dynamical fluctuations; for a bin $\Delta y = 0.25$, fluctuations increase by a factor of two, to $\sim 1.2\%$. This increase in fluctuations on smaller scales is seen in the cosmic microwave background. In the damped scenario, if the time scale for damping is much less than the spatial correlation length, then fluctuations may be essentially constant with the size of the bin.

CERES at the SPS finds dynamical fluctuations of the mean $p_t \leq 3\%$; STAR finds fluctuations $\sim 8\%$ [1]. Regardless of which scenario applies, the condensate model predicts that the fluctuations are concentrated about the mean, and not due to the tails of the distribution in p_t . Thus the experimental signal should not be significantly affected if one cuts on a relatively narrow bin in momentum. STAR included all particles with $.1 < p_t < 1.5$ GeV; in this model the fluctuations are the same for particles with $.3 < p_t < .7$ GeV. Binning in a narrower region in p_t also serves to distinguish it from other models. DCC's [13] or a chiral critical point [18] produce fluctuations at low p_t ; jets give fluctuations at high p_t .

We can also make qualitative predictions about energy loss [5]; in perturbation theory,

the energy loss for a fast particle with transverse momentum p_t is

$$dE/dx \sim \alpha_s m_{\text{Debye}}^2 / \lambda_g \sim \ell_0^2, \quad (4)$$

where $\alpha_s(p_t)$ is the strong coupling constant, m_{Debye} is the Debye mass, and λ_g is the gluon mean free path. In the condensate model, $m_{\text{Debye}}^2 \sim \ell_0^2$, and so vanishes at T_c . This is also seen directly from the lattice data [8]. Assuming that λ_g is finite at T_c gives the dependence on ℓ_0 given above. In QCD, energy loss becomes small as ℓ_0 does.

PHENIX and STAR presented results which appeared to show suppression of particles at high $p_t : 2 \rightarrow 6$ GeV [4]. This is seen even by comparing peripheral to central collisions, as then no model dependent assumptions need to be made. This suggests that at $\sqrt{s} = 130$ A GeV, one is in a regime in which $\ell_0 \neq 0$. At the SPS, no such suppression is observed, which suggests that $\ell_0 \approx 0$ at $\sqrt{s} = 17$ A GeV. (At RHIC, perhaps the largest ℓ_0 is still < 1 , with $\ell_0 \approx 1$ only at LHC?)

The outstanding question is then how these various signals change between $\sqrt{s} = 17$ A GeV and 130 A GeV; is it a sudden transition, or gradual [3]? If the explosive scenario is correct, one might speculate that the transition region is relatively narrow. However, the relationship between temperature and \sqrt{s} may be far from linear. In saturation models [19], for example, $T \sim (\sqrt{s})^{0.2}$, so a relatively narrow region in T corresponds to a broad region in \sqrt{s} .

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